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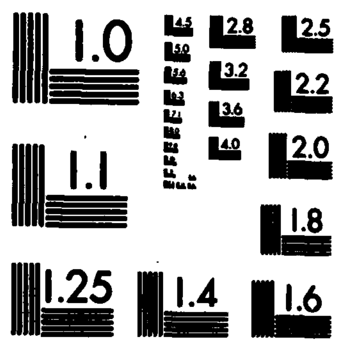

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We develop a set of mode coupling equations which describe the nonlinear evolution of the Rayleigh-Taylor and $E \times B$ gradient drift instabilities which are relevant to the ionosphere. The model is restricted to 2D turbulence in the plane transverse to the magnetic field, and only those modes such that $kL \gg 1$ are considered, where k is the wavenumber and L is the scale length of the density gradient. We show that for a three mode system, the nonlinear equations describing these instabilities correspond exactly to the Lorenz equations which approximately describe the Rayleigh-Benard instability. Following the analysis of Lorenz (1963), it is shown that the three mode system can exhibit a strange attractor with chaotic behavior. Ion inertia plays a critical role in this phenomenon in that if it is neglected (as in the collisional limit), the three mode system does not exhibit chaos and a stable convection pattern results.

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CONTENTS

IONOSPHERIC TURBULENCE: INTERCHANGE INSTABILITIES AND CHAOTIC FLUID BEHAVIOR

I. INTRODUCTION

It is well-known that interchange instabilities produce turbulence in the ionosphere [Ossakow, 1979; Fejer and Kelley, 1980]. The Rayleigh-Taylor instability is believed to cause the intense nighttime equatorial F region turbulence known as equatorial spread F [Ossakow, 1981; Kelley and McClure, 1981]; the $\underset{\sim}{E} \times \underset{\sim}{B}$ gradient drift instability has been invoked to explain high-latitude ionospheric irregularities [Keskinen and Ossakow, 1983], the rapid structuring of barium clouds [Linson and Workman, 1970], and turbulence in equatorial electrojet [Ossakow, 1979]; and the current convective instability may be responsible for turbulence in auroral ionosphere [Ossakow and Chaturvedi, 1979a]. These instabilities are fundamentally similar in nature in that they require a density gradient, and they act to interchange the high and low plasma density regions. However, they have different driving mechanisms: the Rayleigh-Taylor instability is driven by the gravitational force, the $\underset{\sim}{E} \times \underset{\sim}{B}$ gradient drift instability is driven by an ambient electric field or neutral wind, and the current convective instability is driven by a current parallel to the ambient magnetic field. A considerable amount of research has been devoted to the study of these instabilities and their application to ionospheric turbulence. Since, in general, it is the nonlinear phase of the instabilities that is observed, it is important to identify the salient characteristics of this phase (e.g., wave amplitude at saturation, power spectra, etc.).

Interchange instabilities, as they are applied to the ionosphere, can be divided into two categories: collisional and inertial. The collisional limit considers $\nu_{in} >> \omega$ (where ν_{in} is the ion-neutral collision frequency

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and ω is the wave frequency) while the inertial limit considers $\omega >> \nu_{in}$. In the ionosphere, the transition from the collisional limit to the inertial limit typically occurs in the altitude regime ~ 500 km. The bulk of the nonlinear research on interchange instabilities in the ionosphere (both theory and simulation) has been restricted to the collisional domain. The purpose of this paper is to extend the nonlinear theory of interchange instabilities into the inertial regime.

In brief, we develop a set of mode coupling equations which describe the nonlinear evolution of the Rayleigh-Taylor and $\underset{\sim}{E} \times \underset{\sim}{B}$ gradient drift instabilities. The model is restricted to 2D turbulence in the plane transverse to the magnetic field, and only those modes such that $kL >> 1$ are considered, where k is the wavenumber and L is the scale length of the density gradient. We do not consider the current convective instability for simplicity because it requires a component of the wavenumber parallel to $\underset{\sim}{B}$ [Chaturvedi and Ossakow, 1979b]. We show that our mode coupling equations have the same structure as the equations describing the Rayleigh-Benard instability; in particular, when we consider a three mode system, we show that the equations correspond exactly to the Lorenz equations which approximately describe the Rayleigh-Benard instability [Saltzman, 1962; Lorenz, 1963]. Following the analysis of Lorenz (1963), it is shown that the three mode system can exhibit a strange attractor with chaotic behavior [Ruelle and Takens, 1971]. Ion inertia plays a critical role in this phenomenon in that if it is neglected (as in the collisional limit), the three mode system does not exhibit chaos and a stable convection pattern results.

The organization of this letter is as follows. In the next section we present the geometry, assumptions, and derivation of the nonlinear mode coupling equations. In Section III we present our results for a three mode

system and show the correspondence to the Rayleigh-Benard problem. We derive a criterion for the onset of chaotic turbulence involving the dissipation parameters (i.e., diffusion coefficient, ion-neutral collision frequency). Finally, in the last section we summarize our findings and discuss the implications for ionospheric turbulence.

II. DERIVATION OF NONLINEAR EQUATIONS

The plasma configuration and assumptions used in the analysis are described as follows. We consider an ambient magnetic field in the z-direction ($\underset{\sim}{B} = B_0 \hat{e}_z$), a gravitational acceleration in the -x-direction ($\underset{\sim}{g} = -g \hat{e}_x$), a density gradient in the x-direction ($n = n_0(x)$ and $\partial n_0/\partial x > 0$), and an ambient electric field in the y direction ($\underset{\sim}{E} = E_0 \hat{e}_y$). We assume two dimensional perturbations in the x-y plane (transverse to $\underset{\sim}{B}$) such that $\underset{\sim}{k} = k_x \hat{e}_x + k_y \hat{e}_y$ with $k_x L \gg 1$ and $k_y L \gg 1$ where $L = [\partial \ln n/\partial x]^{-1}$ is the density gradient scale length. For simplicity we assume a cold ion plasma ($T_i = 0$). We consider low frequency turbulence in a weakly collisional plasma such that $\partial/\partial t \ll \Omega_i$, $\nu_{in} \ll \Omega_i$, $\nu_{ie} \ll \Omega_i$, and $\nu_{ei} \ll \Omega_e$ where $\Omega_\alpha = eB_0/m_\alpha c$ is the cyclotron frequency of species α, ν_{in} is the ion-neutral collision frequency, ν_{ie} is the ion-electron collision frequency, and ν_{ei} is the electron-ion collision frequency. We neglect electron-neutral collisions since $\nu_{en}/\Omega_e \ll \nu_{in}/\Omega_i$ in the F region. Finally, we consider electrostatic turbulence and assume quasi-neutrality ($n_e \approx n_i$).

The equations used in the analysis are continuity, momentum transfer, and charge conservation:

$$\frac{\partial n_\alpha}{\partial t} + \nabla \cdot (n \underset{\sim}{V}_\alpha) = 0 \tag{1}$$

$$0 = -\frac{e}{m_e}\left(\mathbf{E} + \frac{1}{c}\,\mathbf{V}_e \times \mathbf{B}\right) - \frac{T_e}{m_e}\frac{\nabla n}{n} - \nu_{ei}(\mathbf{V}_e - \mathbf{V}_i) \tag{2}$$

$$\frac{d\mathbf{V}_i}{dt} = \frac{e}{m_i}\left(\mathbf{E} + \frac{1}{c}\,\mathbf{V}_i \times \mathbf{B}\right) - \nu_{ie}(\mathbf{V}_i - \mathbf{V}_e) - \nu_{in}\,\mathbf{V}_i + \mathbf{g} \tag{3}$$

$$\nabla \cdot \mathbf{J} = \nabla \cdot [n(\mathbf{V}_i - \mathbf{V}_e)] = 0 \tag{4}$$

We perturb (1)-(4) about an equilibrium and let $n = n_0 + \tilde{n}$, $\mathbf{E} = \mathbf{E}_0 - \nabla\tilde{\phi}$, and $\mathbf{V}_\alpha = \mathbf{V}_{\alpha 0} + \tilde{\mathbf{V}}_\alpha$. To lowest order in ν_α/Ω_α the equilibrium drifts are given by $\mathbf{V}_{e0} = -(cT_e/eB)(\partial \ell n\, n/\partial x)^{-1}\,\hat{e}_y$ (the electron diamagnetic drift) and $\mathbf{V}_{i0} = [g/\Omega_i + (\nu_{in}/\Omega_i)(cE_0/B)]\hat{e}_y$ (the ion gravitational drift and ion Pedersen drift, respectively). Note we have chosen $V_{0x} + V_{0x} - cE_{0y}/B$. We solve (2) and (3) for $\tilde{V}_e$ and $\tilde{V}_i$, and substitute these values into (1) and (4). We arrive at the coupled set of equations for $\tilde{n}$ and $\tilde{\phi}$:

$$\frac{\partial}{\partial t}\frac{\tilde{n}}{n_0} - \frac{c}{B}\nabla\tilde{\phi} \times \hat{e}_z \cdot \frac{\nabla n_0}{n_0} - D_e\nabla^2\frac{\tilde{n}}{n_0} = \frac{c}{Bn_0}\nabla\tilde{\phi} \times \hat{e}_z \cdot \nabla\tilde{n} \tag{5}$$

and

$$\left[\frac{\mathbf{g}}{\Omega_i} \times \hat{e}_z + \frac{\nu_{in}}{\Omega_i}\frac{c\mathbf{E}}{B}\right] \cdot \nabla\frac{\tilde{n}}{n_0} - \frac{c}{B}\frac{1}{\Omega_i}\left(\frac{\partial}{\partial t} + \hat{e}_z \times \nabla\tilde{\phi} \cdot \nabla + \nu_{in}\right)\nabla^2\tilde{\phi} = 0 \tag{6}$$

where $D_e = \nu_{ei}\rho_e^2$ is the electron diffusion coefficient and we have assumed $\partial/\partial t >> \mathbf{V}_{i0} \cdot \nabla$. It may be seen that (5) and (6) are mathematically the same as the Rayleigh-Benard equations [Eqs. (17) and (18) of Lorenz (1963)] provided the substitution $\nu_{in} \rightarrow \nu\nabla^2$ is made in (6).

III. RESULTS

We now present results of our analysis for a three mode configuration. Prior to this we cast (5) and (6) into dimensionless form. Specifically, we find that (5) and (6) can be written as

$$\frac{\partial \tilde{n}_1}{\partial t_1} = \hat{D}_e \hat{\nabla}^2 \tilde{n}_1 + \hat{\nabla}_y \tilde{\phi}_1 + \hat{\nabla}\tilde{n}_1 \cdot \hat{\nabla}\tilde{\phi}_1 \times \hat{e}_z \tag{7}$$

and

$$\left(\frac{\partial}{\partial t_1} + \hat{e}_z \times \hat{\nabla}\tilde{\phi}_1 \cdot \tilde{\nabla}\right)\hat{\nabla}^2 \tilde{\phi}_1 = - \hat{\nu}_{in}\hat{\nabla}^2 \tilde{\phi}_1 + \hat{\nabla}_y \tilde{n}_1 \tag{8}$$

where $t_1 = \gamma_0 t$, $\hat{\nu}_{in} = \nu_{in}/\gamma_0$, $\tilde{n}_1 = (\tilde{n}/n_0)(L/\lambda)$, $\tilde{\phi}_1 = \tilde{\phi}(c/B\lambda^2\gamma_0)$, $\hat{\nabla} = \lambda\nabla$, $\hat{D}_e = D_e/\lambda^2\gamma_0$, $\gamma_0 = [(g + \nu_{in}cE_0/B)/L]^{1/2}$, and λ is half of the maximum wavelength permitted (i.e., $\lambda < 2L$).

We consider the following perturbations for $\tilde{\phi}_1$ and $\tilde{n}_1$: $\tilde{\phi}_1 = X_1 \sin x \sin y$ and $\tilde{n}_1 = Y_1 \sin x \cos y + Z_1 \sin 2x$ where we have taken $k_x = k_y$ with $k_x = \lambda^{-1}$ and $k_y = \lambda^{-1}$ for simplicity. Here, x and y represent the x and y spatial coordinate normalized to λ, and only the coefficients X_1, Y_1, and Z_1 are assumed to be time dependent. Substituting $\tilde{n}_1$ and $\tilde{\phi}_1$ into (7) and (8) we obtain the following set of coupled ordinary differential equations:

$$\dot{X} = - \sigma X + \sigma Y \tag{9}$$

$$\dot{Y} = - Y + rX - XZ \tag{10}$$

$$\dot{Z} = - 2Z + XY \tag{11}$$

where the dot over a variable indicates a time derivative, $X = X_1 \sqrt{r\sigma}$, $Y = Y_1 \sqrt{r\sigma}/2\hat{\nu}_{in}$, $Z = rZ_1$, $\tau = 2\hat{D}_e t_1$, $\sigma = \hat{\nu}_{in}/2\hat{D}_e$ and $r = (4\hat{D}_e\hat{\nu}_{in})^{-1}$. Equations (9)-(11) correspond exactly to the equations solved by Lorenz (1963) for the Rayleigh-Benard instability (with the exception that Lorenz' parameter b is equal to 2 in our case).

Following Lorenz (1963), (9)-(11) can be analyzed to determine (i) the nonlinear fixed states of the system, and (ii) the stability of these fixed states as a function of r and σ. The fixed states are determined by the condition $\dot{X} = \dot{Y} = \dot{Z} = 0$. If $r < 1$, the only stable steady state is given by $X_0 = Y_0 = Z_0 = 0$. This represents the state of no convection, i.e., the usual equilibrium state upon which linear stability analysis is performed; in fact, $r = 1$ is the point of marginal linear stability of the equilibrium. For $r > 1$, the equilibrium is unstable resulting in convection and the interchange of high and low density regions. However, $r > 1$ allows two additional fixed states, given by $X_0 = Y_0 = \pm [Z(r - 1)]^{1/2}$ and $Z_0 = r - 1$; these correspond to convection cells of either positive or negative vorticity. The linear instability may thus "saturate" by attracting to one of these new fixed states with X_0, Y_0, and Z_0 providing estimates for the saturation amplitudes of $\tilde{n}$ and $\tilde{\phi}$. This is indeed the case for $1 < r < r_c \equiv \sigma(\sigma + 5)/(\sigma - 3)$ or $\sigma < 3$. In this range of r, the fixed states are stable and the orbit in X, Y, Z phase space asymptotes to one of the nontrivial fixed points. An example is shown in Fig. 1 ($\sigma = 10$, $r = 15$) where the projection of the orbit on to the X-Z plane is plotted.

For $\sigma > 3$ and $r > r_c$, the saturated convection pattern described above is itself unstable. No other fixed stable states exist; this means that

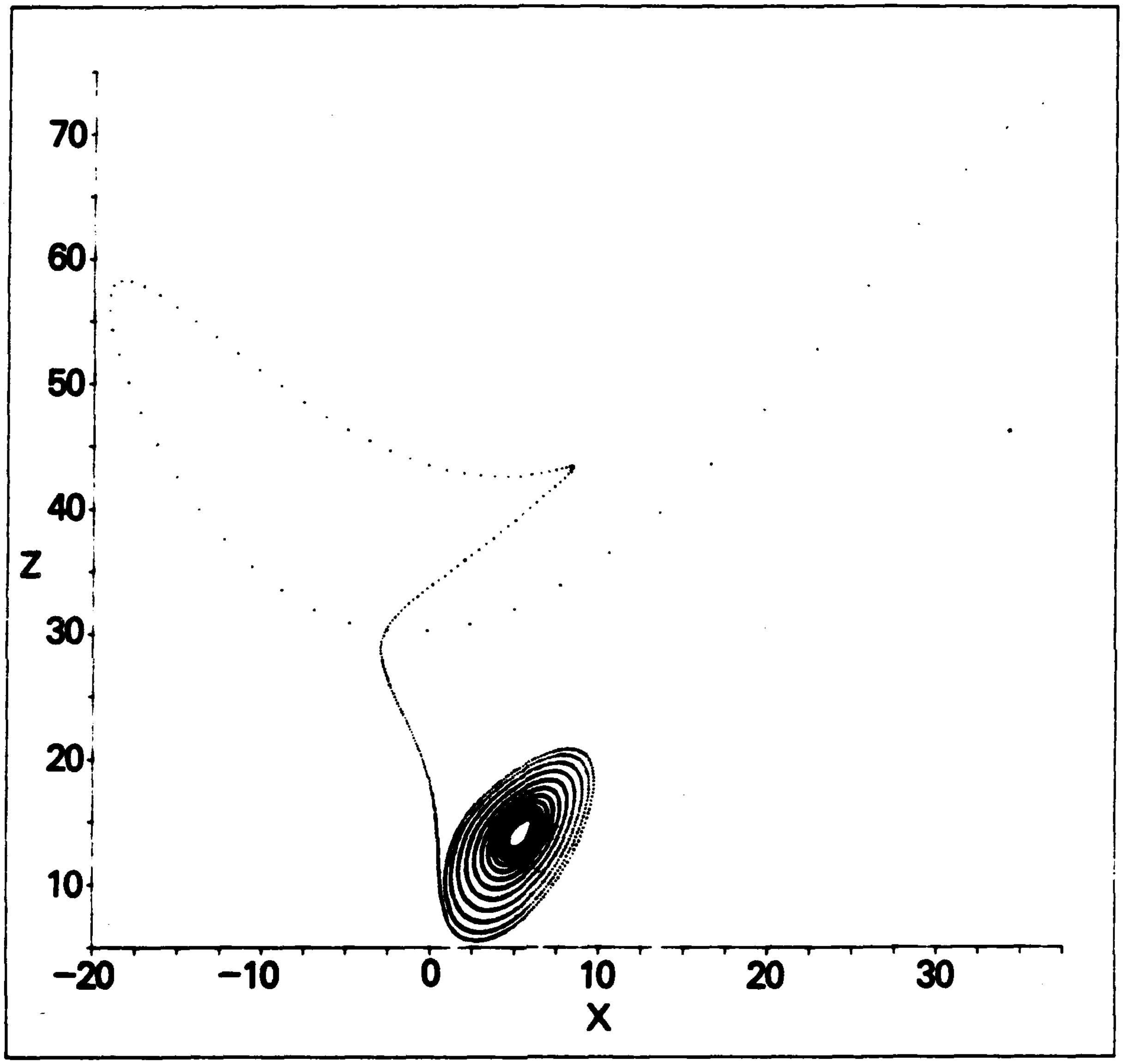

Fig. 1. Depicted is an approach to the non-trivial steady state attractor to (9)-(11) projected onto the X-Z plane. The parameters used for the numerical simulation are $\sigma = 10$ and $r = 15$; note that $r_c = 21.4$ so that $1 < r < r_c$.

the amplitudes X, Y, and Z oscillate in intensity in periodic or chaotic fashion. The magnitude of these oscillations cannot be determined analytically and numerical analysis of (9) - (11) is required. Nevertheless, some general features may be discerned: even though the phase space orbit does not approach a single point, it does lie in a bounded, or "attracting", region of phase space. Figure 2 illustrates such an orbit for $\sigma = 10$ and $r = 30$. This orbit in fact tends to encircle either one or the other of the two non-trivial fixed states. There is, however, no periodicity for the case shown: the transition from encircling one or the other fixed points is seemingly random. In such a case the orbit is "chaotic" and the attracting region is a strange attractor. Since all orbits are unstable, one sees chaotic behavior in the amplitudes of X, Y, and Z [Ruelle and Takens, 1971]. Note that the amplitudes of X and Z in Fig. 2 can fluctuate by more than a factor of 2.

An important point to be recognized concerning the application of this theory to interchange instabilities is the following. As noted earlier, we can consider two limits: collisional and inertial. In the collisional limit $(\nu_{in} \gg \partial/\partial t)$, $\dot{X} = 0$ in (9) so that $X = Y$ and the problem reduces to solving only two coupled differential equations. For this case, the convection states are always <u>stable</u>. This is the situation that has been considered in most previous analytic treatments of ionospheric interchange instabilities [Rognlien and Weinstock, 1974; Chaturvedi and Ossakow, 1977, 1979a,b] although the stability of the convection states was not analyzed. Thus, for this simple three mode system, we find that <u>ion inertia is required for unstable convection patterns to occur</u>. We comment that the nonlinear behavior of Rayleigh-Taylor instability in the inertial limit has been considered by Hudson (1978). However, Hudson (1978) did not

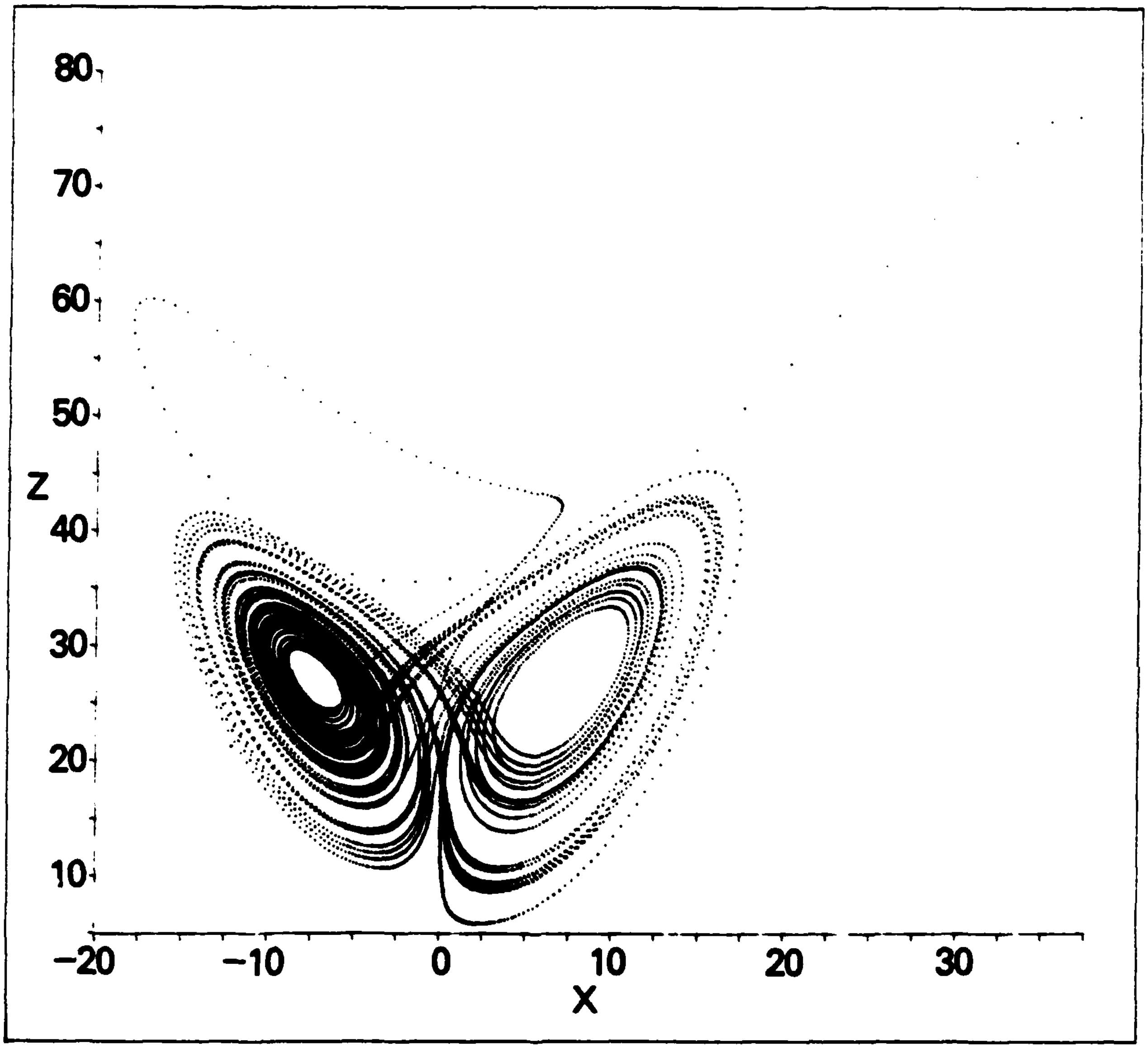

Fig. 2. A strange attractor for (9)-(11) is illustrated as a projection onto the X-Z plane. The parameters used are $\sigma = 10$ and $r = 30$. Note that $r_c = 21.4$ so that $r > r_c$. All periodic orbits are unstable, as well as the three equilibrium points. The attractor lies in a finite volume of space.

find chaotic behavior because it was assumed that $\partial/\partial t = -i\omega$ where ω is the linear eigenfrequency in (6). This ad hoc assumption effectively reduces the problem to two differential equations and leads to results similar to the collisional limit.

IV. CONCLUDING REMARKS

We have shown that interchange instabilities relevant to ionospheric turbulence (Rayleigh-Taylor and $\underset{\sim}{E} \times \underset{\sim}{B}$ gradient drift), can be studied in the context of chaotic attractor theory. In particular, we demonstrate that for a simple system (three modes) the equations governing these instabilities are exactly the same as those that govern the Rayleigh-Benard instability [Saltzman, 1962; Lorenz, 1963]. The analogy between these two instabilities is the following. The Rayleigh-Benard instability is driven by a temperature gradient and convects "hot and cold" fluid elements; the temperature gradient is maintained by a heat source at one end and a heat sink at the other. The Rayleigh-Taylor and $\underset{\sim}{E} \times \underset{\sim}{B}$ gradient drift instabilities are driven by a density gradient and convect "heavy and light" fluid elements; in the case of the ionosphere, the density gradient could be maintained by photoionization at one end and recombination at the other. We have shown that these interchange instabilities can exhibit both stable and unstable convection patterns in this system. A crucial point is that unstable convection only results if inertial effects are important, i.e., $\nu_{in} \lesssim \partial/\partial t$.

For application to the ionosphere, we consider the stability of the fixed states of the Rayleigh-Taylor instability in the equatorial ionosphere, and the $\underset{\sim}{E} \times \underset{\sim}{B}$ gradient drift instability in the high latitude auroral ionosphere. Three observations may be made. First, the nonlinear

fixed states given by X_0, Y_0, and Z_0 correspond to the saturated potential and density fluctuation amplitudes. We note the density fluctuation amplitude associated with Y_0 is $\tilde{n}/n_0 = 4(\rho_e/L)(\nu_{in}\nu_{ei}/\gamma_0^2)^{1/2}$, while the amplitude associated with Z_0 is $\tilde{n}/n_0 = \lambda/L$. These estimates agree with previous results [Chaturvedi and Ossakow, 1977, 1979] and yield density fluctuations of several percent for typical F region parameters. These estimates can vary by more than a factor of 2 when the instabilities are in the "chaotic regime" (see Fig. 2). Second, since $\sigma = \lambda^2\nu_{in}/2\nu_{ei}\rho_e^2$ using the normalizations listed after (8), we find that $10^3 \lesssim \sigma \lesssim 10^5$ for the F region ionosphere (200–800 km) where we have taken $\lambda = 500$ m, $\rho_e = 1.5$ cm, $\nu_{in} = 2.4 \times 10^{11}\ T_e^{1/2} n_n\ \sec^{-1}$ [Strobel and McElroy, 1970], $\nu_{ei} = (\lambda_{ei}/3.5 \times 10^5)(n_e/T_e^{3/2})\ \sec^{-1}$ [Braginskii, 1965; Johnson, 1961] where T_e is the electron temperature in eV, n_n is the neutral gas density, n_e is the electron density, $\lambda_{ei} = 23.4 - 1.15 \log n_e + 3.45 \log T_e$, and the neutral densities were obtained from a Jacchia (1975) model neutral atmosphere. Third, since $\sigma \gg 1$, the critical value of r is given by $r_c \approx \sigma$. For the Rayleigh-Taylor instability, the condition for unstable fixed states can thus be written as $\nu_{in} < (g/2L)^{1/2}$. For $g = 9.8$ m/sec^2 and L = 10 km, we find that unstable behavior can occur for $\nu_{in} < 0.02\ \sec^{-1}$ which corresponds to altitudes greater than 400–500 km in the equatorial F region ionosphere. For the $\underset{\sim}{E} \times \underset{\sim}{B}$ gradient drift instability, the condition for an unstable fixed state is $\nu_{in} < cE_0/2BL$. Taking $cE_0/B = 6 \times 10^2$ m/sec and L = 10^2 km [Weber et al., 1984], we find that chaotic behavior can occur for $\nu_{in} < 6 \times 10^{-3}\ \sec^{-1}$ which corresponds to altitudes greater than roughly 500 km in the high latitude F region.

We comment that the values of r and σ relevant to ionospheric plasmas are considerably different than those used in most mathematical studies

(i.e., similar to those used in Figs. 1 and 2). However, we have also performed calculations for large r and σ and have found similar behavior; namely, the transition from a stable attractor to a strange attractor for sufficiently large r. This is in agreement with previous studies [Fowler and McGuinness, 1982].

Finally, we have only considered a simple three mode system which could be argued is unrealistic. It is known that the nonlinear behavior of the Lorenz equations for more than three modes can be different from the three mode system [Ott, 1981]. For example, we have shown that for three modes, the nonlinear fixed states are always stable in the collisional limit (i.e., $\nu_{in} > \gamma_0$). However, we have also developed a pseudo-spectral code which follows the evolution of a many mode system (~ 120 modes). Preliminary results indicate that chaotic fluid behavior can also occur in the collisional limit, in sharp contrast to the simple three mode result. A report on the nonlinear dynamics of the many mode system will follow shortly.

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